Probing the Universe After Cosmological Recombination Through the Effect of Neutral Lithium on the Microwave Background Anisotropies

Abraham Loeb

Astronomy Department, Harvard University, 60 Garden Street, Cambridge, MA 02138;
aloeb@cfa.harvard.edu

ABSTRACT

I show that neutral lithium couples strongly to the cosmic microwave background (CMB) through its 6708Å resonant transition after it recombines at $z \sim 500$. At observed wavelengths of $\lesssim 335\mu$m, the CMB anisotropies are significantly altered since the optical depth for resonant scattering by neutral lithium is substantial, $\tau_{\rm LiI} \sim 0.5$. The scattering would suppress the original anisotropies, but will generate strong new anisotropies in the CMB temperature and polarization on sub-degree scales ($\ell \gtrsim 100$). Observations at different wavelengths in this spectral regime can probe different thin slices of the early universe. The anisotropy noise contributed by continuum radiation from foreground far–infrared sources could be taken out by subtracting maps at slightly different wavelengths. Detection of the above effects can be used to study structure at $z \lesssim 500$ and to constrain the primordial abundance and recombination history of lithium.

Subject headings: Cosmology: theory — cosmic microwave background

1. Introduction

The dynamical coupling between the baryons and the Cosmic Microwave Background (CMB) during the epoch of cosmological recombination ($z \sim 10^3$) is commonly assumed to be mediated only through Thomson scattering (Peebles & Yu 1970). This coupling becomes ineffective when hydrogen recombines, at around the same time when the universe becomes transparent, and hence it leaves a clear signature on the temperature anisotropies of the CMB (White, Scott, & Silk 1994, and references therein). Forthcoming satellite missions, such as MAP[1] in 2001 or Planck[2] in 2007, offer the sensitivity to probe deviations from this simple model at the percent level (Scott 1999; Yu, Spergel, & Ostriker 2001).

---

[1]http://map.gsfc.nasa.gov

[2]http://astro.estec.esa.nl/Planck

In this *Letter*, I examine the potential significance of the coupling between the cosmic gas and the CMB due to resonant line transitions of neutral atoms. For such transitions, the integral of the absorption cross-section over frequency,

$$\int \sigma(\nu)d\nu = f_{12}\left(\frac{\pi e^2}{m_e c}\right),\tag{1}$$

is typically many orders of magnitude larger than for Thomson scattering, where f_{12} is the absorption oscillator strength (Rybicki & Lightman 1979). For example, the Lyα transition of hydrogen provides an average cross-section which is seven orders of magnitude larger than the Thomson value, when averaged over a frequency band as wide as the resonant frequency itself $[(\Delta\nu/\nu) \sim 1]$. Since the abundance of atoms in excited states is highly suppressed by large Boltzmann factors at the recombination redshift, we consider only transitions from the ground state. For helium, these transitions correspond to photon frequencies that are too far on the Wien tail of the CMB spectrum and only couple to a negligible fraction of the radiation energy density. Hence, we focus our discussion on hydrogen and lithium.

In §2, we derive the resonant–line opacity and drag force for atoms that move relative to a blackbody radiation field. In §3, we consider the cosmological implications of these results for hydrogen and lithium. Finally, §4 summarizes the main conclusions of this work. Throughout the paper, we adopt the density parameters $\Omega_0 \approx 0.3$ for matter and $\Omega_b = 0.04$ for baryons, and a Hubble constant $H_0 \approx 70$ km s^{-1} Mpc^{-1}. The effect of the cosmological constant can be neglected at the redshifts of interest here.

2. Drag Force and Opacity Due to a Resonant Line

Consider a single atom moving at a non-relativistic peculiar velocity $\mathbf{v}$ relative to an isotropic blackbody radiation field of temperature T. Because of its motion, the atom sees a dipole anisotropy in the radiation temperature, $T_a(\theta) = T[1 + (v/c)\cos\theta]$, where θ is the angle relative to its direction of motion (Peebles 1993). We assume that the atom couples to the radiation through a resonant line of frequency ν_0 with $(h\nu_0/kT) \gg 1$. Hence, the line resonates in the Wien tail of the spectrum, for which the radiation energy density per unit frequency (in erg cm^{-3} Hz^{-1}) is given by,

$$u_\nu(\theta) = \frac{8\pi h\nu^3}{c^3}\exp\left\{-\frac{h\nu}{kT_a(\theta)}\right\} \approx \frac{8\pi h\nu^3}{c^3}\exp\left\{-\frac{h\nu}{kT}\right\}\left(1 + \frac{h\nu}{kT}\frac{v}{c}\cos\theta\right),\tag{2}$$

where we have used the fact that $[(h\nu/kT)(v/c)\cos\theta] \ll 1$. Absorption of a resonant photon by the atom is followed by isotropic emission in the atom's rest frame, and so the net drag force (momentum transfer per unit time) is given by,

$$F = -\int_0^\infty d\nu \ \frac{1}{2}\int_{-1}^1 d\cos\theta \ \ \sigma(\nu)u_\nu(\theta)\times\cos\theta \quad .\tag{3}$$

Since the absorption cross-section is sharply peaked around ν_0, we may use equation (1) to get,

$$\mathbf{F} = -\left(\frac{8\pi^2 h^2 f_{12}\nu_0^4}{3c^3 kT}\frac{e^2}{m_e c^2}\exp\left\{-\frac{h\nu_0}{kT}\right\}\right)\mathbf{v} \quad . \tag{4}$$

For an ensemble of atoms of density n_a, which are all moving at the same peculiar velocity $\mathbf{v}$, the force per unit volume is given by $n_a\mathbf{F}$. If the resonant atoms are embedded within a hydrogen plasma of total (free+bound) proton density n_p and are strongly coupled to the plasma through binary particle collisions, then the cosmological equation of motion for the gas as a whole reads,

$$(m_p n_p + m_a n_a)\left(\frac{d\mathbf{v}}{dt} + H\mathbf{v}\right) = n_a\mathbf{F}, \tag{5}$$

or equivalently,

$$\frac{dv}{dt} + Hv = -\frac{v}{t_a}, \tag{6}$$

where $H = (\dot{a}/a)$ is the Hubble expansion rate, and t_a is the characteristic time over which the peculiar velocity of the gas is damped due to the drag force on the resonant atoms,

$$t_a(t) \equiv \frac{3m_p\lambda_0^4}{8\pi^2 h^2 c f_{12}}\frac{m_e c^2}{e^2}\frac{kT}{X_a}\exp\left\{\frac{h\nu_0}{kT}\right\}. \tag{7}$$

Here, $\lambda_0 = c/\nu_0$ is the resonant wavelength, m_p is the proton mass, and $X_a \equiv n_a/(n_p + A_a n_a)$, where $A_a = m_a/m_p$ is the atomic weight of the resonant atoms. In the cosmological context, $T = 2.725 \text{ K} \times (1+z)$ is the CMB temperature (Mather et al. 1999). Both T and X_a are functions of cosmic time, and so equation (6) admits the solution, $v = [v_0/a(t)]\exp\{-\int dt/t_a(t)\}$, where $v_0 = const$ and $a(t)$ is the cosmological scale factor. In general, the right-hand-side of equation (6) should include other terms which source the peculiar velocity field, such as gravitational or pressure forces due to density inhomogeneities. The significance of the drag force can be calibrated in terms of the Hubble time, $H^{-1}(z)$, through the product Ht_a. For the redshifts of interest, $H(z) \approx \sqrt{\Omega_0}H_0(1+z)^{3/2}$.

The assumption of an isotropic radiation background is only valid if the moving fluid element is *optically-thin* to resonant scattering. In the opposite regime of high opacity, the background photons are isotropized locally through multiple scatterings and exert a reduced drag force on an atom embedded deep inside the moving fluid element. The Sobolev optical depth for resonant scatterings in an expanding and nearly uniform universe is given by (Sobolev 1946; Dell'Antonio & Rybicki 1993),

$$\tau_a(z) = f_{12}\frac{\pi e^2}{m_e c}\frac{\lambda_0 n_a(z)}{H(z)} = \frac{A_{21}(g_2/g_1)\lambda_0^3 n_a(z)}{8\pi H(z)} \quad , \tag{8}$$

where A_{21} is the transition probability per unit time for spontaneous emission between the two energy levels (in s^{-1}), and (g_2/g_1) is the ratio between the statistical weights of the excited and ground states. The optical depth is dominated by the velocity gradient of the Hubble flow, because

gravitationally–induced peculiar velocities perturb this flow only slightly at the early cosmic times of interest here.

For $\tau_a \gg 1$, the flux of resonant photons inside the moving fluid element is reduced by a factor of $\sim \tau_a$ (Rybicki & Hummer 1978), and so the drag force in equation (3) is lowered by the same factor. This reduction has dramatic consequences for the resonant Lyα transition of neutral hydrogen, as we show next.

3. Implications For Cosmological Recombination

3.1. Hydrogen

For the $1S$–$2P$ Lyα transition of hydrogen, $\lambda_0 = 1216$Å , $A_{21} = 6.262 \times 10^8$ s^{-1}, $f_{12} = 0.4162$, and $(g_2/g_1) = 3$, yielding an exceedingly large Sobolev optical–depth, $\tau_H = 7.5 \times 10^8 X_H[(1+z)/10^3]^{3/2}$, where X_H denotes the neutral fraction of hydrogen. The thermally–broadened Lyα line has an effective cross–section of $\langle\sigma\rangle \sim 3 \times 10^{-17}$ cm^2 (Peebles 1993) and a corresponding mean-free-path of $1/n_H\langle\sigma\rangle \sim 10^{14}$ cm, at $z \sim 10^3$. Only on smaller scales can Lyα viscosity damp velocity gradients effectively. On the much larger scales where the CMB anisotropies are measurable, the drag time in equation (7) is lengthened by a factor of $\sim \tau_H$, which makes it very much longer than the Hubble time. The situation is similar for deuterium. Balmer line transitions of hydrogen have a small optical depth, but the occupation probability of their lower level is suppressed by an exceedingly large Boltzmann factor, $\sim \exp\{-4.35 \times 10^4/(1+z)\}$.

The resulting drag force is small, even when including the enhancement in the Lyα flux beyond the blackbody value, due to the Lyα photons released during recombination (~ 1 Lyα photon per hydrogen atom; see Dell'Antonio & Rybicki 1993 for the distorted spectrum).

3.2. Lithium

3.2.1. Drag Force and Opacity

Lithium has a closed $n = 1$ shell with two electrons, and one electron outside this shell. The transition between the ground state ($2S$) and the first excited state ($2P$) has $\lambda_0 = 6708$Å and $A_{21} = 3.69 \times 10^7$ s^{-1}, with $f_{12} = 0.247$ and $(g_2/g_1) = 1$ for the 2^2S–$2^2P^0_{1/2}$ transition, and $f_{12} = 0.494$, $(g_2/g_1) = 2$ for the 2^2S–$2^2P^0_{3/2}$ transition (Radzig & Smirnov 1985; Yan & Drake 1995). We therefore get from equations (7) and (8),

$$Ht_{\mathrm{LiI}} = 4.96 \times 10^{-4} \left(\frac{1+z}{10^3}\right)^{5/2} \left(\frac{X_{\mathrm{LiI}}}{3.8 \times 10^{-10}}\right)^{-1} \exp\left\{\frac{7.89 \times 10^3}{1+z}\right\}, \tag{9}$$

and

$$\tau_{\mathrm{LiI}} = 2.82 \left(\frac{X_{\mathrm{LiI}}}{3.8 \times 10^{-10}}\right) \left(\frac{1+z}{10^3}\right)^{3/2} . \tag{10}$$

Figure 1 depicts the redshift dependence of the drag time in equation (9) and compares it to the drag time due to Thomson scattering in the recombining plasma,

$$H t_{\mathrm{e}} = \frac{3 m_p c H}{4 \sigma_T a T^4 X_e} = 4.92 \times 10^{-3} X_{\mathrm{e}}^{-1} \left(\frac{1+z}{10^3}\right)^{-5/2} , \tag{11}$$

where $X_{\mathrm{e}} = (n_{\mathrm{e}}/n_p) = n_{\mathrm{e}}/(n_{\mathrm{e}} + n_{\mathrm{H}})$ is the ionization fraction of hydrogen. We calculated $X_{\mathrm{e}}(z)$ from the updated version of the standard recombination model (Seager, Sasselov, & Scott 2000). Despite the low abundance of lithium, its resonant drag force is not much lower than the Thomson drag force at $z \sim 500$, when roughly half of the lithium atoms are expected to have recombined (Palla, Galli, & Silk 1995; Stancil, Lepp, & Dalgarno 1996, 1998).

The net effect of the lithium drag is highly sensitive to the recombination history of lithium and hydrogen. Previous discussions (e.g., Stancil et al. 1996, 1998) did not include the delaying effect of Lyα photons on the recombination history of hydrogen (Peebles 1993), and therefore underestimated the free electron fraction and correspondingly the neutral lithium abundance at $z \lesssim 10^3$. Another process with an opposite sign that was omitted, involves the excitation of neutral lithium from its ground level by the CMB radiation field. For this effect, it is important to take account of the Lyman and Balmer series distortions of the CMB spectrum (Dell'Antonio & Rybicki 1993), which redshift into resonance at lithium recombination. In fact, the lithium chemistry in the early universe might provide a way for inferring the existence of these distortions. Finally, the recombination history of ^{7}Be should also be followed since a substantial fraction of the ^{7}Li forms through electron-capture in the ^{7}Be produced by Big-Bang nucleosynthesis (T. Walker, private communication). Since the half-life of neutral ^{7}Be is only 53 days, this conversion can be regarded as instanteneous on the cosmological timescale, as soon as ^{7}Be recombines. Careful calculations of the neutral lithium fraction as a function of redshift are required in order to quantify the imprint of the lithium drag on the CMB anisotropies (Dalgarno, Loeb, & Stancil 2001).

The lithium chemistry could also be affected by a large drift velocity between the Li I fluid and the hydrogen plasma. To examine this effect, we consider the equation of motion for the Li I fluid (Burgers 1969),

$$m_{\mathrm{LiI}} \left(\frac{d\mathbf{v}_{\mathrm{LiI}}}{dt} + H\mathbf{v}_{\mathrm{LiI}}\right) = -m_{\mathrm{LiI}} \frac{\mathbf{v}_{\mathrm{LiI}}}{X_{\mathrm{LiI}} t_{\mathrm{LiI}}} + m_{\mathrm{p}} W_{\mathrm{LiI}} \left(\mathbf{v} - \mathbf{v}_{\mathrm{LiI}}\right) . \tag{12}$$

Here $W_{\mathrm{LiI}} = \Sigma_k \mu_{\mathrm{k:Li}} W_{\mathrm{k:Li}}$ is the collision rate of a Li atom with the gas, where $W_{\mathrm{k:LiI}} = n_k \langle \sigma_{\mathrm{k:LiI}} v \rangle$, and $\mu_{\mathrm{k:Li}} = [m_k m_{\mathrm{Li}}/m_{\mathrm{p}}(m_k + m_{\mathrm{Li}})]$. The index k runs over all other particle species in the gas which are assumed to move at a common bulk velocity $\mathbf{v}$ because they are not subject to the resonant drag force. The geometric cross-section of Li I atoms, $\sigma_{\mathrm{LiI}} \sim 10^{-15}$ cm^2, couples them to the rest of the gas on a time scale $\sim 10^7$ s, much shorter than the Hubble time, $H^{-1} \sim 3 \times 10^{13}$ s

at $z \sim 10^3$. Hence, we may ignore the terms on the left-hand-side in equation (12), as well as the perturbed gravitational force which is of the same order. The remaining terms give the solution,

$$\frac{(v - v_{\mathrm{LiI}})}{v} = \frac{1}{1 + Y} \tag{13}$$

where $Y = W_{\mathrm{LiI}} t_{\mathrm{LiI}} X_{\mathrm{LiI}} / A_{\mathrm{Li}}$. Substituting $A_{\mathrm{Li}} = 7$, we find that $W_{\mathrm{LiI}} \approx 10^{-7}[(1 + z)/10^3]^{3.5}$ s^{-1}, and that $Y \gg 1$ at $z \lesssim 400$, after lithium recombination. At higher redshifts, the assumption that $v_{\mathrm{LiI}} = v$ seems not to apply and one may wonder whether the drag time in equation (6) should be multiplied by a factor of $(v/v_{\mathrm{LiI}}) = (1 + Y)/Y > 1$, since the drag force acts on $\mathbf{v}_{\mathrm{LiI}}$ and not on $\mathbf{v}$. In this context it is crucial to realize that the Li I nuclei do not maintain their identity as neutral atoms; in fact, they continuously form and get ionized on a timescale much shorter than a Hubble time. The cross-section for Coulomb collisions of Li$^+$ is far greater than that of Li I, $\sigma_{\mathrm{Li}^+} = 2\sqrt{\pi} e^4 \ln \Lambda / (kT)^2 = 2.7 \times 10^{-11}[(1 + z)/10^3]^{-2}$ cm^2, where $\ln \Lambda \approx 20$ is the Coulomb logarithm. The fraction of time spent by a Li nucleus in an ionized state is the ionization fraction, $f_+ = X_{\mathrm{Li}^+}/(X_{\mathrm{LiI}} + X_{\mathrm{Li}^+})$, and so σ_{LiI} should be replaced by $[f_+ \sigma_{\mathrm{Li}^+} + (1 - f_+)\sigma_{\mathrm{LiI}}]$ and the drag force term should be multiplied by $(1 - f_+)$ in equation (12). We consequently find that the substantial ionization fraction of the lithium fluid (Stancil et al. 1996, 1998) diminishes its drift velocity relative to the gas at all relevant redshifts, $z \lesssim 10^3$. Hence, non-thermal relative velocities may be ignored in the Li I chemistry.

At $z \lesssim 10^3$, the mean-free-path for resonant scattering by lithium exceeds the relevant spatial scales for the CMB anisotropies (10^{23}–10^{24} cm), even if the lithium abundance is taken to be somewhat higher than the value, $X_{\mathrm{SBBN}} = 3.8 \times 10^{-10}$, predicted by the latest deuterium measurements and Standard Big-Bang Nucleosynthesis (Burles, Nollett, & Turner 2001; see also Walker et al. 1991; Smith, Kawano, & Malaney 1993). Note that values as high as $X_{\mathrm{LiI}} \sim 10^{-8}$ were suggested by models of inhomogeneous Big-Bang nucleosynthesis (Applegate & Hogan 1985; Sale & Mathews 1986; Mathews et al. 1990), and would result in $H t_{\mathrm{LiI}} \lesssim 10$ for velocity gradients on sufficiently small scales with $\tau_{\mathrm{LiI}} \lesssim 1$.

3.2.2.   Effects on CMB Anisotropies

If a substantial fraction of lithium recombines by $z \sim 500$ (as predicted by Palla et al. 1995 or Stancil et al. 1996, 1998), then the flux of the original anisotropies will be suppressed by the absorption factor, $\exp(-\tau_{\mathrm{LiI}})$, at observed wavelengths below (6708Å $\times 500$) = 335μm, for which $\tau_{\mathrm{LiI}} \approx 0.6(X_{\mathrm{LiI}}/X_{\mathrm{SBBN}})$. The resonant scattering effect is different from that caused by Thomson scattering at the epoch of reionization (Hu & White 1997), in that it has a larger optical depth and it occurs within a much thinner shell of gas (for a sufficiently narrow band of observed wavelengths). Consequently, it should induce new first-order anisotropies [multiplied by $(1 - \exp(-\tau_{\mathrm{LiI}})]$ due to the coherence of the velocity field in the thin scattering shell. The dominant contribution to the new anisotropies would come from the Doppler effect. In contrast to reionization, there should be no severe cancellations of the Doppler effects from line-of-sight velocity fluctuations at

large wavenumbers, because the last scattering surface (or "visibility function") for the resonance is extremely thin. The increase in amplitude of sub-horizon fluctuations between $z = 10^3$ and $z = 500$ would enhance the new anisotropies relative to the original ones and introduce a distortion of order unity in the original anisotropy spectrum on these scales. The anisotropy spectrum would be modified at multipole moments $\ell \gtrsim 100$, around and below the angular scale of the first acoustic peak (which is a factor of $\sim \sqrt{6} = 2.4$ smaller than the horizon scale at $z \sim 500$), where the original anisotropy amplitude reaches its maximum ($\ell \approx 220$). Resonant scattering would also result in enhanced polarization anisotropies on the same angular scales (Chandrasekhar 1960). Since the scattering is done by lithium, the anisotropies will reflect the peculiar velocity of the lithium fluid which, as shown above, is not expected to deviate significantly from the velocity of the gas as a whole.

Since the temperature fluctuations are in the Wien tail of the CMB spectrum, they translate to brightness fluctuations (in erg s^{-1}cm^{-2}sr^{-1}Hz^{-1}) of a much larger contrast,

$$\frac{\Delta B_\nu}{B_\nu} = \left(\frac{d\ln B_\nu}{d\ln T}\right)\frac{\Delta T}{T} = \left(\frac{h\nu}{kT}\right)\frac{\Delta T}{T} = 15.78 \times \left(\frac{500}{1+z}\right)\frac{\Delta T}{T} \quad , \tag{14}$$

where z is the redshift being probed and we have substituted $B_\nu(T) \propto \exp(-h\nu/kT)$. Hence the first acoustic peak would correspond to brightness fluctuations of up to $\sim 5 \times 10^{-4}$.

At a given observed wavelength, the redshift thickness of the scattering shell is only determined by two contributions: (i) the possibility that a photon at this wavelength was either absorbed by the 2^2S–$2^2P^0_{3/2}$ transition with $\lambda_0 = 6707.76$Å or at a slightly lower redshift by the 2^2S–$2^2P^0_{1/2}$ transition with $\lambda_0 = 6707.91$Å (Radzig & Smirnov 1985), yielding $[\Delta z/(1+z)] = 2.2 \times 10^{-5}$; and (ii) the thermal width of the line[3], $[\Delta z/(1+z)] = (\Delta\lambda/\lambda)_{\rm th} \approx 10^{-5}$. In practice, the minimum thickness would be limited by the band–width of the detector. The maximum band–width and minimum angular resolution should be tuned to match the expected coherence length of the peculiar velocity field at $z \sim 500$, which corresponds to $(\Delta\lambda/\lambda) \sim 0.2$ and $\Delta\theta \sim 6'$ or equivalently $\ell \sim 10^3$.

Observations at different far–infrared (FIR) wavelengths would probe different thin slices of structure in the early universe. The anisotropy pattern would vary gradually as a function of wavelength and "3D tomography" is in principle possible, by which one may probe correlations in the structure along the line-of-sight. The feasibility of this measurement might be severely compromised by the anisotropies of the FIR foreground. Unfortunately, 350μm is just the wavelength where the CMB intensity blends into the extragalactic FIR foreground (see Fig. 1 in Scott et al. 2001). However, the contribution of discrete sources (see Fig. 6 in Knox et al. 2001) could be separated out through observations at shorter wavelengths with a higher angular resolution. Furthermore, continuum foreground emission would result in similar anisotropies for all

---

[3]Since $\tau_{\rm LiI} \lesssim 1$, the photon does not need to redshift away from the line center by more than the line width in order to escape to infinity.

wavelengths within a band width $(\Delta\lambda/\lambda) \sim 1$, while the CMB anistropies would vary considerably across such a band since its width corresponds to $\Delta z/(1+z) \sim 1$. Subtraction of anisotropy maps at slightly different wavelengths within such a band can therefore be used to isolate out the CMB component (as long as these wavelengths do not overlap with strong emission lines from the foreground sources). For $\ell \lesssim 100$, the anisotropies should have the same power-spectrum as measured by MAP or Planck at much longer wavelengths. The predictable correlations between the temperature and polarization patterns of the CMB and their general statistical properties could be used to further enhance the signal-to-noise ratio (Zaldarriaga & Loeb 2001). Obviously, these difficult observations need to be made through holes in the Galactic dust emission at high Galactic latitudes.

4. Conclusions

The 6708Å resonant transition of neutral lithium provides a previously unexplored coupling between the baryons and the CMB after cosmological recombination. The lithium opacity is substantial at observed wavelengths $\lesssim 335[(1+z_{\mathrm{LiI}})/500]\mu$m, where z_{LiI} is the redshift at which $\sim 50\%$ of the lithium recombines. At these wavelengths, the CMB anisotropies on sub-degree angular scales would be significantly altered, due to the finite optical depth for resonant scattering, $\tau_{\mathrm{LiI}} \approx 0.5[(1+z_{\mathrm{LiI}})/500]^{3/2}$. This scattering would generate new temperature and polarization anisotropies at mutipole moments $\ell \gtrsim 100$. The detection of these anisotropies depends on the prospects for cleaning the expected anisotropies of the FIR foreground (Knox et al. 2001), and may be feasible through the subtraction of anisotropy maps for pairs of slightly different wavelengths.

The relevant wavelength range overlaps with the highest frequency channel of the Planck mission (352μm) and with the proposed baloon–borne Explorer of Diffuse Galactic Emissions[4] (EDGE) which will survey 1% of the sky in 10 wavelength bands between 230–2000μm with a resolution ranging from $6'$ to $14'$ (see Table 1 in Knox et al. 2001). However, a new instrument design with multiple narrow bands $(\Delta\lambda/\lambda \lesssim 0.1)$ at various wavelengths in the range $\lambda = 250$–350μm and with angular resolution $\sim 6'$, is necessary in order to optimize the detection of the lithium signature on the CMB anisotropies.

The effects discussed in this *Letter* depend critically on the recombination history of lithium, and emphasize the need to explore the related lithium chemistry in more detail in the future.

---

[4]http://topweb.gsfc.nasa.gov

REFERENCES

Applegate, J. H. & Hogan, C. J. 1985, Phys. Rev. D, 31, 3037

Burgers, J. M. 1969, Flow Equations For Composite Gases (New York: Academic)

Burles, S., Nollett, K. M., & Turner, M. S. 2001, ApJL, in press; astro-ph/0010171

Chandrasekhar, S. 1960, Radiative Transfer, (New-York: Dover), pp. 50-53

Dalgarno, A., Loeb, A., & Stancil, P. 2001, in preparation

Dell'Antonio, I. P., & Rybicki, G. B. 1993, in ASP Conf. Ser. 51, Observational Cosmology, , ed. G. Chincarini et al. (San Francisco: ASP), 548

Knox, L, Cooray, A., Eisenstein, D., & Haiman, Z. 2001, ApJ, 550, 7

Mather, J. C., Fixsen, D. J., Shafer, R. A., Mosier, C., & Wilkinson, D. T. 1999, ApJ, 512, 511

Mathews, G. J., Alcock, C. R., & Fuller, G. M. 1990, ApJ, 349, 449

Hu, W., & White, M. 1997, ApJ, 479, 568

Palla, F., Galli, D., & Silk, J. 1995, ApJ, 451, 44

Peebles, P. J. E. 1993, Principles of Physical Cosmology, (Princeton: Princeton Univ. Press), pp. 154-157, 168–176

Peebles, P. J. E. & Yu, J. T. 1970, ApJ, 162, 815

Radzig, A. A., & Smirnov, B. M. 1985, Reference Data on Atoms, Molecules, and Ions (New York: Springer), p. 224

Rybicki, G. B., & Hummer, D. G. 1978, ApJ, 219, 654

Rybicki, G. B., & Lightman, A. P. 1979, Radiative Processes in Astrophysics, (New York: Wiley), p. 102

Sale, K. E. & Mathews, G. J. 1986, ApJL, 309, L1

Seager, S., Sasselov, D. D., & Scott, D. 2000, ApJS, 128, 407

Scott, D. 1999, to appear in "Beyond the Desert '99", edited by H.V. Klapdor-Kleingrothaus; astro-ph/9911325

Scott, D. et al. 2001, to appear in UMass/INAOE conference proceedings on Deep Millimeter Surveys, eds. J. Lowenthal and D. Hughes; astro-ph/0010337

Smith, M. S., Kawano, L. H., & Malaney, R. A. 1993, ApJS, 85, 219

Sobolev, V. V. 1946, Moving atmospheres of Stars (Leningrad: Leningrad State University [in Russian]; English translation: 1960 (Cambridge: Harvard Univ. Press)

Stancil, P. C., Lepp, S., & Dalgarno, A. 1996, ApJ, 458, 401

————————————————. 1998, ApJ, 509, 1

Yan, Z. & Drake, G. W. F. 1995, Phys. Rev. A, 52, 4316

Yu, Q., Spergel, D. N., & Ostriker, J. P. 2001, ApJ, submitted, astro-ph/0103149

Walker, T. P., Steigman, G., Kang, H., Schramm, D. M., & Olive, K. A. 1991, ApJ, 376, 51

White, M., Scott, D., & Silk, J. 1994, ARA&A, 32, 319

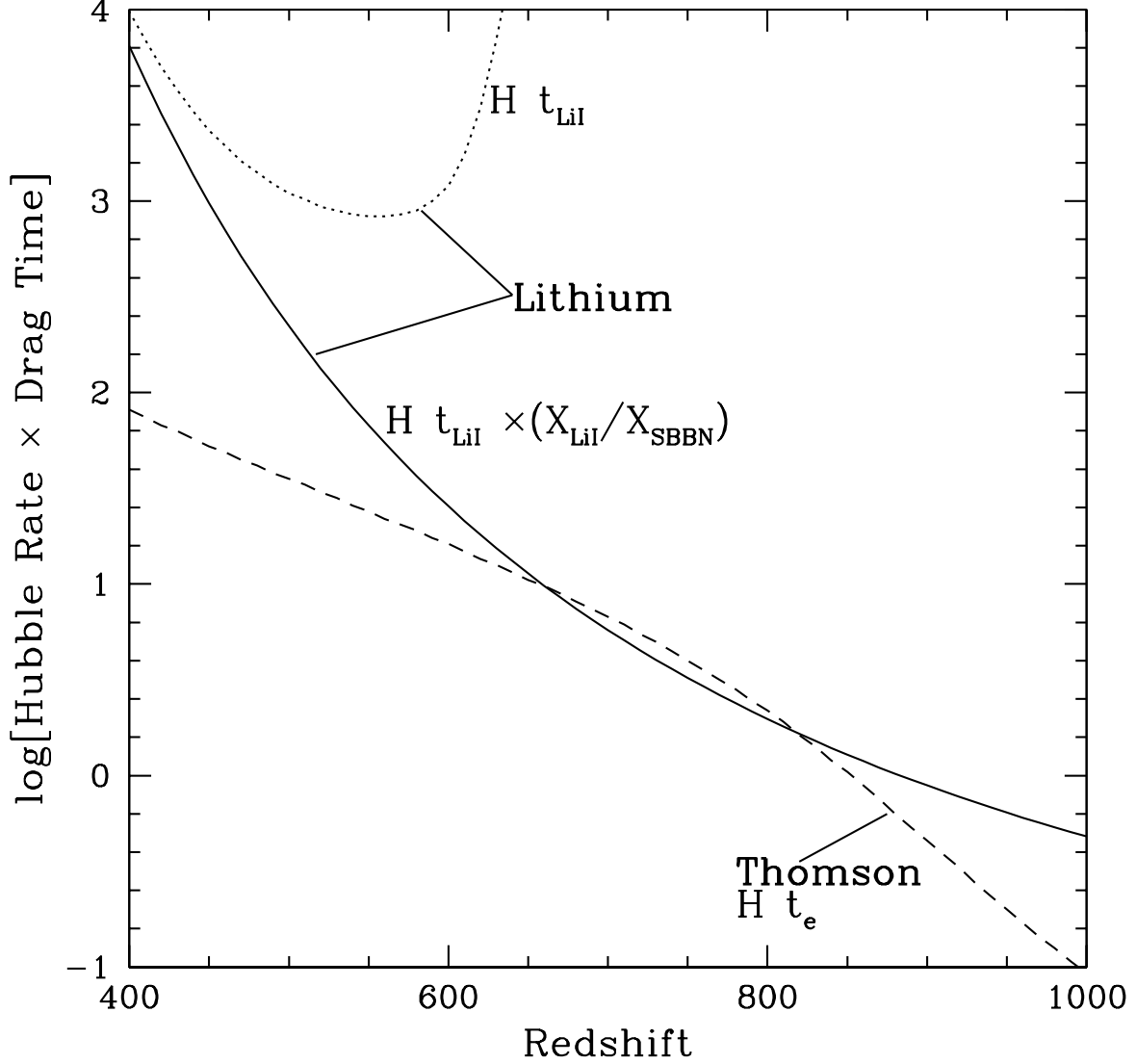Zaldarriaga, M., & Loeb, A. 2001, in preparation

Fig. 1.— Redshift evolution of the dynamical drag time due to the Li 6708Å resonance with the microwave background. The drag time shown by the solid line should be divided by the neutral fraction of the primordial lithium abundance ($X_{\mathrm{SBBN}} = 3.8 \times 10^{-10}$, see Burles et al. 2001), as a function of redshift. The dotted line shows the result for one of the recombination histories of lithium in the literature (Palla et al. 1995). For comparison, the Thomson drag time is plotted as the dashed line, taking account of the full recombination history of hydrogen. In all cases, the drag times are normalized by the Hubble time, $H^{-1}(z)$.